

The 750 GeV diphoton excess and its explanation in 2-Higgs Doublet Models with a real inert scalar multiplet

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Abstract

We discuss a possible explanation of the recently observed diphoton excess at around 750 GeV as seen by the ATLAS and CMS experiments at the Large Hadron Collider. We calculate the cross section of the diphoton signature in 2-Higgs Doublet Models with the addition of a real isospin scalar multiplet without a vacuum expectation value, where a neutral component of such a representation can be a dark matter candidate. We find that the branching fraction of an additional CP-even Higgs boson H from the doublet fields into the diphoton mode can be significantly enhanced, by up to a factor of 10^3 , with respect to the case of the simple 2-Higgs Doublet Model. Such a sizable enhancement is realized due to multi-charged inert particle loops entering the $H \rightarrow \gamma\gamma$ decay mode. Through this enhancement, we obtain a suitable cross section of the $gg \rightarrow H \rightarrow \gamma\gamma$ process to explain the data with the fixed mass of H being 750 GeV.

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I. INTRODUCTION

Recently, there was a quite exciting report about a new excess at around 750 GeV in the diphoton invariant mass distribution at the 13 TeV run of the CERN Large Hadron Collider (LHC) from both ATLAS and CMS groups. The excess was observed with a local significance around 3.6σ at ATLAS [1] using 3.2 fb^{-1} data and around 2.6σ at CMS [2] using 2.6 fb^{-1} data. Most instinctively, this excess can be explained in such a way that there is a new neutral spin-0 resonant particle¹ with a mass of about 750 GeV which decays into two photons. If this conjecture is correct and the excess is confirmed by further data, this will represent direct evidence for the existence of new physics Beyond the Standard Model (BSM).

So far, there appeared a number of papers to describe the excess since it was reported. For example, in Refs. [5], the excess was explained via a scalar boson resonance by introducing an effective Lagrangian with dimension five operators. It was also explained in a model with an extra isospin singlet scalar field [6–8], that with an extra doublet scalar field including the minimal supersymmetric Standard Model (SM) [8–10] and that with an extra triplet scalar field [11, 12]. A connection between the Dark Matter (DM) physics and the diphoton excess has been discussed in Ref. [13].

We discuss the possibility to explain this excess in 2-Higgs Doublet Models (2HDMs) supplemented by a suitable additional scalar representation. The latter addition to a standard 2HDM spectrum is needed, as Ref. [9] has already pointed out that the gluon fusion process $gg \rightarrow H/A \rightarrow \gamma\gamma$ within 2HDMs, where H and A are the additional CP-even and CP-odd Higgs bosons, respectively, is not sufficient to produce the event yield required, essentially because the branching fraction of $H/A \rightarrow \gamma\gamma$ is typically of $\mathcal{O}(10^{-3})$ smaller than the necessary value.

In fact, there are two ways to reproduce the required cross section, of order 10 fb: to have an enhancement in (i) the gluon fusion cross section and/or (ii) in the diphoton branching fraction. In this paper, we consider the second possibility, by introducing an additional real inert scalar multiplet without a non-zero Vacuum Expectation Value (VEV) in 2HDMs with a softly-broken Z_2 symmetry. Such an inert scalar multiplet is for example motivated in the “minimal DM scenario” discussed in Refs. [14, 15], where a neutral component of the multiplet can be a viable DM candidate. Thanks to the introduction of the inert multiplet, the branching fraction of $H \rightarrow \gamma\gamma$ is

¹ Although there could exist in principle a spin 2 particle as an alternative explanation, it seems difficult to construct a phenomenologically reasonable model to include such a state. In contrast, a spin 1 resonant state cannot directly decay into the diphoton final state [3], but it can decay into the triphoton state via a scalar boson inter mediation [4].

	u_R	d_R	e_R	Q_L, L_L	ξ_u	ξ_d	ξ_e
Type-I	-	-	-	+	$\cot \beta$	$\cot \beta$	$\cot \beta$
Type-II	-	+	+	+	$\cot \beta$	$-\tan \beta$	$-\tan \beta$
Type-X	-	-	+	+	$\cot \beta$	$\cot \beta$	$-\tan \beta$
Type-Y	-	+	-	+	$\cot \beta$	$-\tan \beta$	$\cot \beta$

TABLE I: The Z_2 charge assignments for the SM fermion fields. The ξ_f factors are given for each 2HDM type.

significantly enhanced, indeed by a factor of 10^2 - 10^3 depending on the isospin of the inert multiplet.

The plan of the paper is as follows. In Sec. II, we define our model. In Sec. III, we calculate the gluon fusion production cross section and the decay branching ratio into the diphoton state of the additional neutral Higgs bosons. By combining these, we show the cross section of the diphoton process. The cutoff scale of our model is estimated by using one-loop renormalization group equations (RGEs). We then conclude in Sec. IV. In Appendix, we present the set of beta functions for all the relevant coupling constants at the one-loop level.

II. THE MODEL

We consider an extension of 2HDMs where the scalar sector is composed of two isospin doublets Φ_1 and Φ_2 with hypercharge $Y = 1/2$ and a real ($Y = 0$) inert scalar multiplet χ with isospin T with a null VEV. In particular, we consider the case for $T = 2, 3$ and 4, or equivalently χ is assumed to be an isospin quintuplet, septet and nonet, respectively. We note that the maximal allowed value of T has been obtained to be 4 from perturbative unitarity arguments² [16].

The active sector involving Φ_1 and Φ_2 which have non-zero VEVs is similar to that of 2HDMs, so that there appears flavor changing neutral currents (FCNCs) at the tree level. In order to forbid such phenomena, we impose a softly-broken discrete Z_2 symmetry [18] under which the scalar fields transform as $(\Phi_1, \Phi_2, \chi) \rightarrow (+\Phi_1, -\Phi_2, +\chi)$. Under the Z_2 symmetry, there are four types of Yukawa interactions [19, 20] and they are called Type-I, Type-II, Type-X and Type-Y [21] depending on the Z_2 charge assignment for right-handed fermions. In Table I, the Z_2 charge assignment for all fermion fields is shown in each types of Yukawa interactions.

² In Ref. [17], the scale dependence on dimensionless scalar couplings has been calculated in a model with one $SU(2)_L$ doublet Higgs field and a higher isospin multiplet using one-loop renormalization group equations. It has been clarified that if we have a $T \geq 3/2$ multiplet, the Landau Pole appears below the Planck scale.

The most general scalar potential under the $SU(2)_L \times U(1)_Y \times Z_2$ invariance with CP-conservation is given by [22]

$$\begin{aligned}
V(\Phi_1, \Phi_2, \chi) = & \mu_1^2 \Phi_1^\dagger \Phi_1 + \mu_2^2 \Phi_2^\dagger \Phi_2 - \mu_3^2 (\Phi_1^\dagger \Phi_2 + \text{h.c.}) + \mu_\chi^2 \chi^\dagger \chi \\
& + \frac{1}{2} \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \frac{1}{2} \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 |\Phi_1^\dagger \Phi_2|^2 + \frac{1}{2} \lambda_5 [(\Phi_1^\dagger \Phi_2)^2 + \text{h.c.}] \\
& + \lambda_\chi (\chi^\dagger \chi)^2 + \lambda'_\chi (\chi^\dagger T^a T^b \chi)^2 + \rho_1 (\Phi_1^\dagger \Phi_1) (\chi^\dagger \chi) + \rho_2 (\Phi_2^\dagger \Phi_2) (\chi^\dagger \chi),
\end{aligned} \tag{1}$$

where T^a ($a = 1, 2, 3$) are the $(2T + 1)$ dimensional representation of the $SU(2)$ generator. We note that the operator $(\chi^\dagger T^a \chi)$ is identically zero [22]. We also note that for the case of $T = 2$, the λ'_χ term is written by the λ_χ term, so that we can take $\lambda'_\chi = 0$. On the other hand for the $T = 3$ and 4 case, the λ'_χ term gives the independent combination of the component fields of χ .

There are five physical scalar states from Φ_1 and Φ_2 as in 2HDMs, i.e., one pair of singly-charged Higgs bosons H^\pm , a CP-odd Higgs boson A and two CP-even Higgs bosons h and H (with $m_h < m_H$), where h is assumed to be the discovered Higgs boson with a mass of about 125 GeV. The ratio of the two doublet VEVs v_1 and v_2 is defined as $\tan \beta = v_2/v_1$ and the two VEVs satisfy $v_1^2 + v_2^2 = v^2 = (\sqrt{2}G_F)^{-1/2} \simeq (246 \text{ GeV})^2$. The mass formulae for these Higgs bosons are the same as those in 2HDMs (see e.g., Ref. [23] for the explicit expressions).

The squared mass of χ is given by

$$m_\chi^2 = \mu_\chi^2 + \frac{v^2}{2} (\rho_1 \cos^2 \beta + \rho_2 \sin^2 \beta), \tag{2}$$

where all the masses of the component scalar fields in χ are degenerate at the tree level. A non-zero mass splitting can be generated via one-loop corrections whose amount is typically only $\mathcal{O}(100)$ MeV as shown in Ref. [14]. In the following discussion, we thus neglect such a small mass difference.

III. DIPHOTON EXCESS

We now discuss how we can reproduce the diphoton excess at around 750 GeV at the LHC in our BSM scenario. Herein, the additional neutral Higgs bosons H and A can contribute to this excess via the gluon fusion process, i.e., $gg \rightarrow H/A \rightarrow \gamma\gamma$, by taking the mass of H and A to be 750 GeV. The production cross section of $gg \rightarrow H/A$ is calculated as follows

$$\sigma(gg \rightarrow H/A) = \sigma(gg \rightarrow h_{\text{SM}}) \times \frac{\Gamma(H/A \rightarrow gg)}{\Gamma(h_{\text{SM}} \rightarrow gg)}, \tag{3}$$

where h_{SM} is the Higgs boson in the SM with a mass of 750 GeV. The SM cross section $\sigma(gg \rightarrow h_{\text{SM}})$ with $m_{h_{\text{SM}}} = 750 \text{ GeV}$ at the collision energy of 13 TeV is given by about 736 fb, as quoted from

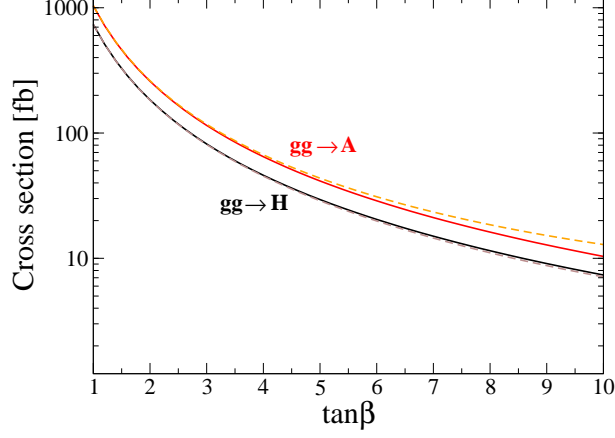


FIG. 1: Cross section of $gg \rightarrow H$ and $gg \rightarrow A$ as a function of $\tan\beta$ for the case of $\sin(\beta - \alpha) = 1$ and $m_A = m_H = 750$ GeV. The solid (dashed) curves show the cases with Type-I and Type-X (Type-II and Type-Y) Yukawa interaction.

the LHC Higgs Cross Section Working Group page [24]. The decay rate is given by

$$\Gamma(\mathcal{H} \rightarrow gg) = \frac{\sqrt{2}G_F\alpha_s^2 m_{\mathcal{H}}^3}{128\pi^3} \left| \sum_f \xi_{\mathcal{H}}^f F_{1/2}^{\mathcal{H}}(m_f) \right|^2, \quad \mathcal{H} = H, h, \quad (4)$$

$$\Gamma(A \rightarrow gg) = \frac{\sqrt{2}G_F\alpha_s^2 m_A^3}{128\pi^3} \left| \sum_f 2T_f^3 \xi_f F_{1/2}^A(m_f) \right|^2, \quad (5)$$

where T_f^3 is the third component of the isospin of f , i.e., $T_f^3 = +1/2$ ($-1/2$) for $f = u$ (d, e), and the $\xi_{\mathcal{H}}^f$ factors are expressed in terms of the ξ_f expressions given in Tab. I by

$$\xi_h^f = \sin(\beta - \alpha) + \xi_f \cos(\beta - \alpha), \quad (6)$$

$$\xi_H^f = \cos(\beta - \alpha) - \xi_f \sin(\beta - \alpha). \quad (7)$$

We note that the decay rate for h_{SM} is obtained from $\Gamma(h \rightarrow gg)$ by taking $\sin(\beta - \alpha) = 1$.

The loop functions for the CP-even Higgs bosons $F_{1/2}^{\mathcal{H}}(m_f)$ and the CP-odd Higgs boson $F_{1/2}^A(m_f)$ are given by

$$F_{1/2}^{\mathcal{H}}(m_f) = -\frac{4m_f^2}{m_{\mathcal{H}}^2} \left[2 - m_{\mathcal{H}}^2 \left(1 - \frac{4m_f^2}{m_{\mathcal{H}}^2} \right) C_0(0, 0, m_{\mathcal{H}}^2, m_f, m_f, m_f) \right], \quad (8)$$

$$F_{1/2}^A(m_f) = -4m_f^2 C_0(0, 0, m_A^2, m_f, m_f, m_f), \quad (9)$$

where C_0 is the usual three-point Passarino-Veltman function [25].

Now, we are ready to calculate the gluon-fusion production cross sections of H and A . In Fig. 1, we show these as a function of $\tan\beta$ in the case of $\sin(\beta - \alpha) = 1$ and $m_H = m_A = 750$

GeV in all the types of Yukawa interactions. The cross section is monotonically reduced as $\tan \beta$ increases, because the top loop contribution is suppressed due to the $\cot \beta$ factor in the $Ht\bar{t}$ and $At\bar{t}$ couplings. The dependence upon the type of Yukawa interaction is not so important in this region of $\tan \beta$. In the case of $\tan \beta < 1$, the top Yukawa coupling becomes too strong to guarantee the validity of a perturbative calculation. Besides, such a parameter region is strongly constrained from the B physics measurements such as B^0 - \bar{B}^0 mixing [26, 27]. We thus do not consider the region with $\tan \beta < 1$ in this paper.

Next, let us discuss the decay rates of the neutral Higgs bosons H and A into the diphoton mode. They are given by

$$\Gamma(\mathcal{H} \rightarrow \gamma\gamma) = \frac{\sqrt{2}G_F\alpha_{\text{em}}^2 m_{\mathcal{H}}^3}{256\pi^3} \times \left| \xi_{\mathcal{H}}^V F_1^{\mathcal{H}}(m_W) + \sum_f Q_f^2 N_c^f \xi_{\mathcal{H}}^f F_{1/2}^{\mathcal{H}}(m_f) - \frac{\lambda_{\mathcal{H}H^+H^-}}{v} F_0^{\mathcal{H}}(m_{H^\pm}) - Q_\chi^2 \frac{\lambda_{\mathcal{H}\chi^\dagger\chi}}{v} F_0^{\mathcal{H}}(m_{\chi^\pm}) \right|^2, \quad (10)$$

$$\Gamma(A \rightarrow \gamma\gamma) = \frac{\sqrt{2}G_F\alpha_{\text{em}}^2 m_A^3}{256\pi^3} \left| \sum_f N_c^f \xi_f Q_f^2 F_{1/2}^A(m_f) \right|^2, \quad (11)$$

where N_c^f is the color factor, Q_f is the electric charge for a fermion f and Q_χ^2 represents the squared sum of the electric charges of the charged component field of χ defined as

$$Q_\chi^2 = T^2 + (T-1)^2 + \cdots + 1^2. \quad (12)$$

The $\xi_{\mathcal{H}}^V$ factor denotes the $\mathcal{H}VV$ ($V = W, Z$) couplings divided by the corresponding $h_{\text{SM}}VV$ couplings, namely, $\xi_h^V = \sin(\beta - \alpha)$ and $\xi_H^V = \cos(\beta - \alpha)$. For the charged scalar loops, the scalar trilinear couplings $\lambda_{\mathcal{H}H^+H^-}$ and $\lambda_{\mathcal{H}\chi^\dagger\chi}$ are defined as the coefficient of the scalar trilinear vertices in the Lagrangian. They are expressed as³

$$\lambda_{hH^+H^-} = \frac{1}{v} [(2M^2 - 2m_{H^\pm}^2 - m_h^2)s_{\beta-\alpha} + 2(M^2 - m_h^2)\cot 2\beta c_{\beta-\alpha}], \quad (13)$$

$$\lambda_{HH^+H^-} = -\frac{1}{v} [2(M^2 - m_H^2)\cot 2\beta s_{\beta-\alpha} + (2m_{H^\pm}^2 + m_H^2 - 2M^2)c_{\beta-\alpha}], \quad (14)$$

$$\lambda_{h\chi^\dagger\chi} = -2v[s_{\beta-\alpha}(\rho_1 c_\beta^2 + \rho_2 s_\beta^2) - c_{\beta-\alpha}s_\beta c_\beta(\rho_1 - \rho_2)], \quad (15)$$

$$\lambda_{H\chi^\dagger\chi} = -2v[c_{\beta-\alpha}(\rho_1 c_\beta^2 + \rho_2 s_\beta^2) + s_{\beta-\alpha}s_\beta c_\beta(\rho_1 - \rho_2)], \quad (16)$$

where M^2 denotes the soft-breaking scale of the Z_2 symmetry defined by $M^2 = \mu_3^2/(\sin \beta \cos \beta)$ [28].

³ Here we use the short-hand notations $s_\beta = \sin \beta$, $c_\beta = \cos \beta$, $s_{\beta-\alpha} = \sin(\beta - \alpha)$ and $c_{\beta-\alpha} = \cos(\beta - \alpha)$.

The loop functions in Eq. (10) are given by

$$F_0^{\mathcal{H}}(m_{\phi^\pm}) = \frac{2v^2}{m_{\mathcal{H}}^2} [1 + 2m_{\phi^\pm}^2 C_0(0, 0, m_{\mathcal{H}}^2, m_{\phi^\pm}, m_{\phi^\pm}, m_{\phi^\pm})], \quad (17)$$

$$F_1^{\mathcal{H}}(m_W) = \frac{2m_W^2}{m_{\mathcal{H}}^2} \left[6 + \frac{m_{\mathcal{H}}^2}{m_W^2} + (12m_W^2 - 6m_{\mathcal{H}}^2) C_0(0, 0, m_{\mathcal{H}}^2, m_W, m_W, m_W) \right]. \quad (18)$$

Notice here that the charged scalar loop contributions to the decay rate only appear for the case of CP-even Higgs bosons, so that the decay rate of A into diphoton is not enhanced by charged scalar loops. Typically, the branching fraction of $A \rightarrow \gamma\gamma$ is of $\mathcal{O}(10^{-5})$ when $\tan\beta \sim 1$ in all types of Yukawa interactions. By looking at Fig. 1, the production cross section of $gg \rightarrow A$ process is about 1 pb at $\tan\beta \sim 1$, so that we obtain that $\sigma(gg \rightarrow A \rightarrow \gamma\gamma) \simeq \sigma(gg \rightarrow A) \times \text{BR}(A \rightarrow \gamma\gamma) \simeq 0.01$ fb. This value is $\mathcal{O}(10^3)$ times smaller than the required cross section to explain the diphoton excess. For this reason, the CP-odd Higgs boson contribution does not help to explain the diphoton excess, and in the following we concentrate on the H contribution, i.e., $gg \rightarrow H \rightarrow \gamma\gamma$.

We now show numerical results for the decay rates of the H state, see Fig. 2. In order to obtain the maximal value of the gluon-fusion cross section, we take $\tan\beta = 1$. In that case, the dependence upon the type of Yukawa interaction only appears in the sign of the bottom quark and tau lepton Yukawa couplings $Hb\bar{b}$ and $H\tau^+\tau^-$ (see Tab. I), respectively, and such a difference can be safely neglected. We thus do not specify the type of Yukawa interaction in the following analysis any further. In addition, we take the SM-like limit $\sin(\beta - \alpha) = 1$ to keep the SM-like Higgs boson h coupling to the gauge bosons hVV and fermions $hf\bar{f}$ to be the same values as those in the SM itself. Such a situation is favored by the LHC Run-I data [29, 30]. Therefore, our benchmark scenario to explain the diphoton excess is defined by the following parameter choices:

$$m_H = m_A = m_{H^\pm} = M = 750 \text{ GeV}, \quad \tan\beta = 1, \quad \sin(\beta - \alpha) = 1 \quad \text{and} \quad \rho_2 = -\rho_1. \quad (19)$$

In the above configuration, we obtain $\lambda_{h\chi^\dagger\chi} = 0$ but $\lambda_{H\chi^\dagger\chi} \neq 0$, so that the χ loop effect to the diphoton decay rate only appears in H . We thus can realize the scenario where the branching fraction of $h \rightarrow \gamma\gamma$ is kept to be the SM value, while that of $H \rightarrow \gamma\gamma$ gets a significant contribution from the χ loop. We note that the H^\pm and W boson loop effects to the $H \rightarrow \gamma\gamma$ decay vanish in this configuration because $\lambda_{HH^+H^-} = \xi_H^V = 0$.

Before we move on the numerical evaluation for the branching fraction, we briefly mention about constraints on the parameter space from theoretical and experimental sources. First, there are bounds from the perturbative unitarity and the vacuum stability which are given by theoretical requirements. The former and latter constraints have been discussed in Ref. [31] and in Ref. [32] in 2HDMs, respectively. In the setup given in Eq. (19), we obtain $\lambda_1 = \lambda_2 = \lambda_3 = m_h^2/v^2 \simeq 0.26$

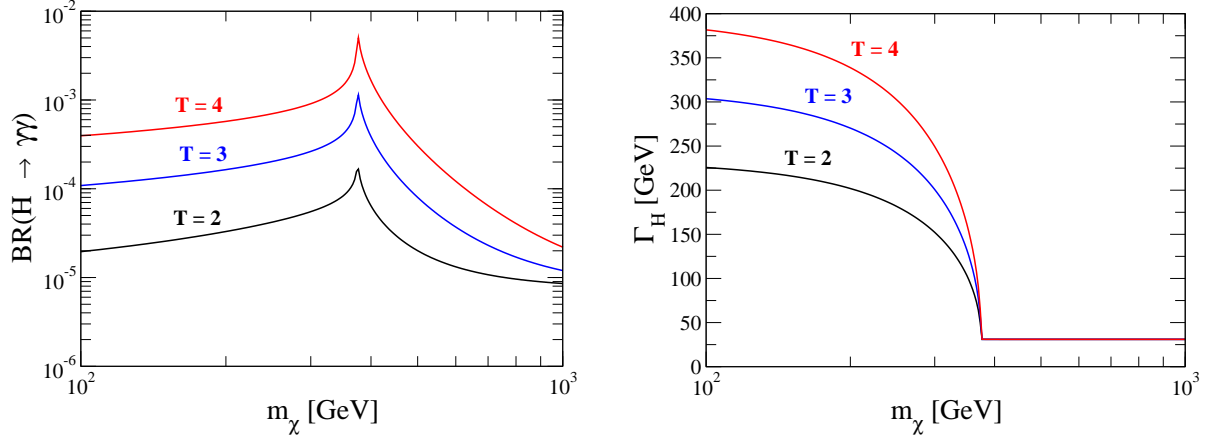


FIG. 2: (Left) branching fraction of $H \rightarrow \gamma\gamma$ as a function of m_χ in the case of $\tan\beta = 1$, $\sin(\beta - \alpha) = 1$ and $m_A = m_{H^\pm} = m_H = M = 750$ GeV. We fix $\rho_1 = -2\sqrt{\pi}$ and $\rho_2 = -\rho_1$. (Right) The total width of H as a function of m_χ for the same parameter set. In both frames, the black, blue and red curves show the cases with $T = 2, 3$ and 4 , respectively.

and $\lambda_4 = \lambda_5 = 0$. This satisfy both the unitarity and vacuum stability constraint in the 2HDM. Of course, we cannot simply apply these constraints to our model because of the introduction of the χ field, i.e., we need to take into account the additional dimensionless parameters ρ_1 , ρ_2 , λ_χ and λ'_χ in the potential. In this paper, we simply argue the perturbativity of these parameters, where the magnitude of these couplings do not exceed a certain value such as $2\sqrt{\pi}$ or 4π .

As the experimental constraints, we first consider the electroweak precision valuables such as S , T and U parameters [33]. The additional contributions to S and T parameters are exactly zero under the setup in Eq. (19) at the one-loop level. Besides, if we neglect the small mass difference between W and Z bosons, the new contribution to the U parameter is also exactly canceled. Regarding to the flavour constraint, we can apply the same bound in the 2HDM to our model up to the leading order calculation because χ does not directly couple to the SM fermions. According to the recent study for the flavour constraints in the 2HDM [27], our benchmark point is allowed. In Ref. [6], the 95% CL upper limit on the cross section for $pp \rightarrow \phi^0 \rightarrow XY$ (ϕ^0 being a new neutral scalar boson with the mass of 750 GeV, and X and Y being the SM particles) is shown at the LHC with the collision energy of 8 TeV, where the gluon fusion production cross section at 8 TeV is about 5 times smaller than that at 13 TeV. We check that in our benchmark scenario, all the predictions of the cross section are below the upper bound⁴ given in [6].

⁴ Only for the $\gamma\gamma$ mode, our prediction can be compatible to the upper limit on the cross section ~ 1.5 fb [6] in the

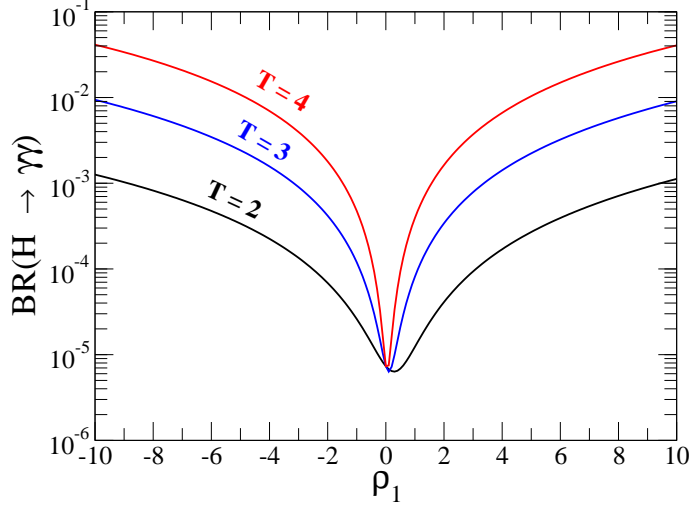


FIG. 3: Branching fraction of $H \rightarrow \gamma\gamma$ as a function of ρ_1 ($= -\rho_2$) in the case of $\tan\beta = 1$, $\sin(\beta - \alpha) = 1$ and $m_A = m_{H^\pm} = m_H = M = 750$ GeV. The black, blue and red curves show the cases with $T = 2, 3$ and 4 , respectively. The mass of χ is fixed to be 375 GeV.

In Fig. 2, we show the m_χ dependence of the branching fraction of $H \rightarrow \gamma\gamma$ (left panel) and the total width Γ_H of H (right panel). We can see in the left figure that the branching fraction becomes the maximum value at around $m_\chi \simeq m_H/2 \simeq 375$ GeV, because of the threshold effect of the χ loop contribution. We note that the H^\pm loop contribution vanishes in the benchmark point given in Eq. (19) because of $\lambda_{HH^+H^-} = 0$. By looking at the right figure, we see the drastic difference in Γ_H between the cases of $m_\chi < m_H/2$ and $m_\chi > m_H/2$. If we take $m_\chi < m_H/2$, the tree level decay mode $H \rightarrow \chi^\dagger\chi$ opens, where $\chi^\dagger\chi$ includes all the possible combination of the component fields in χ ($\chi^\dagger\chi = \chi^0\chi^0, \chi^+\chi^-, \dots$). Conversely, in the case of $m_\chi > m_H/2$, the tree level decay is kinematically forbidden and Γ_H does not depend strongly upon either the isospin T or the mass m_χ and is about 30 GeV, indeed a value compatible with the one fitted to the 750 GeV excess.

In Fig. 3, we show the ρ_1 dependence of the branching fraction of the $H \rightarrow \gamma\gamma$ mode. We take $m_\chi = 375$ GeV to extract the maximal value of the branching fraction for each fixed value of ρ_1 . We note that a slightly larger value of the branching fraction is obtained when we take negative values of ρ_1 as compared to the case of positive values with the same magnitude. The reason is that the contributions between the top loop and the χ loop becomes constructive (destructive) when we take $\rho_1 < 0$ ($\rho_1 > 0$). We find that, in order to obtain the branching fraction to be the

case where the cross section of the diphoton excess, i.e., ~ 5 -10 fb at 13 TeV can be explained.

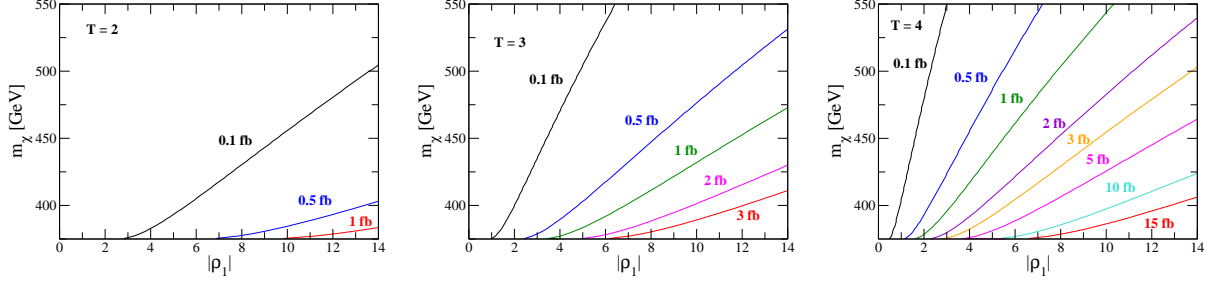


FIG. 4: Contour plots for the cross section of $gg \rightarrow H \rightarrow \gamma\gamma$ on the $|\rho_1|$ - m_χ plane with $\rho_1 < 0$ and $\rho_2 = -\rho_1$. We take $\tan\beta = 1$, $\sin(\beta - \alpha) = 1$ and $m_A = m_{H^\pm} = m_H = M = 750$ GeV. The left, center and right panels show the cases with $T = 2, 3$ and 4 , respectively.

order of 10^{-2} , we need $|\rho_1| \gtrsim 10$ (4) for the case of $T = 3$ (4).

In Fig. 4, we show the contour plots for the cross section of the $gg \rightarrow H \rightarrow \gamma\gamma$ process, where we use the narrow width approximation. The results for $T = 2, 3$ and 4 are shown in the left, center and right panel, respectively. We find that in the case of $\rho_1 = -2\sqrt{\pi}$ and $m_\chi = 375$ GeV, the cross section is given to be about 0.1 fb, 1 fb and 3 fb for $T = 2, 3$ and 4 , respectively.

Finally, we discuss the cutoff scale Λ in our model, wherein running coupling constants become infinity, namely, we have a Landau pole at Λ and our model should be replaced by a fundamental theory which is expected to describe physics from Λ to the Planck scale. In our model, we expect the Landau pole to appear well below the Planck scale, because we introduced rather large coupling constants, ρ_1 and ρ_2 , plus a high isospin scalar multiplet χ . In order to know the cutoff scale, we calculate the running coupling constants by solving the RGEs at the one-loop level. In Appendix, we give the one-loop beta functions for all the coupling constants. We take the initial scale to be $m_H = 750$ GeV and the set of initial values of coupling constants as follows

$$\begin{aligned} \lambda_1(m_H) = \lambda_2(m_H) = \lambda_3(m_H) = \frac{m_h^2}{v^2}, \quad \lambda_4(m_H) = \lambda_5(m_H) = 0, \quad y_t(m_H) = \frac{\sqrt{2}m_t}{v}, \\ \rho_2(m_H) = -\rho_1(m_H), \quad \lambda_\chi(m_H) = \lambda'_\chi(m_H) = 0, \end{aligned} \quad (20)$$

with $m_h = 125$ GeV and $m_t = 173.21$ GeV. Instead of the extraction of the scale $\lambda_i(\Lambda) = \infty$ (λ_i are the coupling constants in our model), we define the cutoff scale⁵ to be the scale giving $\lambda_i(\Lambda) = 4\pi$.

In Fig. 5, we show the cutoff scale Λ as a function of ρ_1 for each case of $T = 2, 3$ and 4 . We see that the model with a larger isospin multiplet has Λ at lower energy scale as it is expected.

⁵ The choice of this critical value, taking 4π , 8π and so on, is not so important, because once one of the coupling constants exceeds $\mathcal{O}(10)$, then such a coupling quite rapidly blows up.

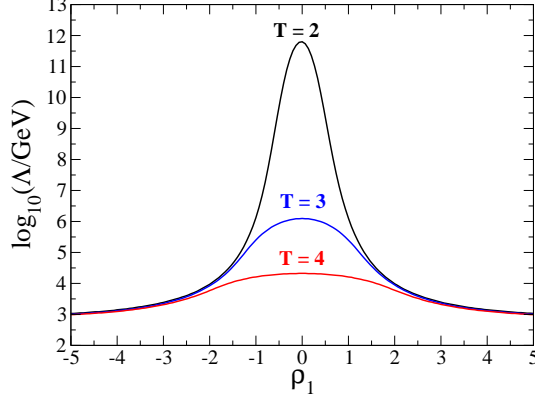


FIG. 5: The cutoff scale Λ as a function of ρ_1 for the case of $T = 2$ (black), $T = 3$ (blue) and $T = 4$ (red). We use the set of initial values of the running coupling constants given in Eq. (20).

However, once $|\rho_1| \gtrsim 2$ is taken, Λ is given to be below the order 10 TeV scale in all the cases. Therefore, if we require $\Lambda > 1$ TeV, then the maximally allowed cross section of the diphoton process is about less than 0.1 fb, 0.5 fb and 1 fb in the case of $T = 2, 3$ and 4, respectively.

Before closing this section, we would like to briefly comment on the possibility of one-loop induced $H \rightarrow VV$ ($V = W, Z$) decays. Although in the above analysis we take $\sin(\beta - \alpha) = 1$ in order to eliminate the tree level $H \rightarrow VV$ decays, these could appear at the one-loop level, where the top quark runs in both the $H \rightarrow W^+W^-$ and $H \rightarrow ZZ$ diagrams, while χ does so only in the $H \rightarrow W^+W^-$ diagram (there are no $\chi^\dagger\chi Z$ and $\chi^\dagger\chi ZZ$ couplings due to the real scalar nature of χ). Since χ has a large isospin charge and a large coupling to the H state, we expect the χ loop to give a significant contribution to the $H \rightarrow W^+W^-$ decay, which can reach the upper limit on the cross section of $pp \rightarrow H \rightarrow W^+W^-$. The upper limit was derived to be 40 fb with 95% CL from the LHC Run-I data [6]. The one-loop induced decay rate of $H \rightarrow W^+W^-$ is expressed as

$$\begin{aligned} \Gamma(H \rightarrow W^+W^-) &= \left(\frac{g^2}{16\pi^2}\right)^2 \frac{1}{16\pi m_H} |c_{\text{eff}}^{HWW}|^2 \left(3 - \frac{m_H^2}{m_W^2} + \frac{m_H^4}{4m_W^4}\right) \sqrt{1 - \frac{4m_W^2}{m_H^2}} \\ &\simeq (0.19 \text{ GeV}) \times \frac{|c_{\text{eff}}^{HWW}|^2}{m_H^2}, \end{aligned} \quad (21)$$

where c_{eff}^{HWW} is the dimensionful one-loop induced HW^+W^- coupling factorized by $g^2/(16\pi^2)$. Furthermore, the cross section of the $gg \rightarrow H$ process is about 160 fb [24] at 8 TeV, which sets a 95% CL upper bound on the branching fraction of $H \rightarrow W^+W^-$ mode at about 0.25. Using the total width of the H state to be 30 GeV, we obtain the upper bound on $|c_{\text{eff}}^{HWW}|$ to be about 4.7 TeV. Suppose that the effective coupling c_{eff}^{HWW} is naively given by $\sum_{Q=0}^T [T(T+1) - Q]^2 \lambda_{H\chi^* \chi^* Q} v^2 C_0(m_W^2, m_W^2, m_H^2; m_\chi, m_\chi, m_\chi)$ then we get $|c_{\text{eff}}^{HWW}|$ to be less than 2.9 TeV, 8.7 TeV

and 19 TeV for the case of $T = 2, 3$ and 4, with $\rho_1 = 1$ ($\rho_2 = -\rho_1$). Therefore, our prediction of the branching fraction of $H \rightarrow W^+W^-$ could be comparable or even above the upper bound. However, in order to get a more precise result, we would need to perform the renormalization prescription of the HW^+W^- vertex as well as take into account the destructive top loop effect, which is beyond the scope of the current paper.

IV. CONCLUSIONS

We have discussed the extension of the softly-broken Z_2 symmetric 2HDMs, where the scalar sector is composed of two active complex doublets and a real inert scalar multiplet χ with an isospin $T(= 2, 3, 4)$ whose lightest neutral component could be a DM candidate. In this model, we have investigated the possibility to explain the diphoton excess at around 750 GeV recently observed by the LHC experiments with 13 TeV energy. In our model, the additional active neutral Higgs boson H and A contribute to the diphoton process via $gg \rightarrow H/A \rightarrow \gamma\gamma$. In order to explain this excess, we need an $\mathcal{O}(10^3)$ enhancement in the branching fractions of $H/A \rightarrow \gamma\gamma$ as compared to the standard 2HDMs. Such an enhancement can be realized by the loop effects induced by charged inert particles. We have shown that the CP-odd Higgs boson A contribution does not help to explain the excess because the branching fraction of $A \rightarrow \gamma\gamma$ cannot be enhanced by the χ loop diagrams. In contrast, the branching fraction of $H \rightarrow \gamma\gamma$ can be significantly enhanced by such a χ loop effect. Upon using the narrow width approximation, we have found that, if we take the coupling ρ_1 to be $-2\sqrt{\pi}$, the cross section $gg \rightarrow H \rightarrow \gamma\gamma$ can be about 0.1 fb, 1 fb and 3 fb in the case of $T = 2, 3$ and 4, respectively. When we allow rather strong couplings such as $\rho_1 = -6$, we then obtain a cross section of about 0.5 fb, 3 fb and 10 fb in the case of $T = 2, 3$ and 4, respectively. We have checked where the cutoff scale Λ appears in our model. By solving the RGEs at the one-loop level, we obtained Λ to be about 10^{12} , 10^6 and $10^{4.5}$ GeV for $T = 2, 3$ and 4 when $\rho_1 \simeq 0$. This scale becomes smaller than $\mathcal{O}(10)$ TeV when we take $|\rho_1| \gtrsim 2$ for $T = 2, 3$ and 4, and, if we require $\Lambda > 10$ TeV, then the maximally allowed cross section of the diphoton process is less than about 0.1 fb, 0.5 fb and 1 fb in the case of $T = 2, 3$ and 4, respectively.

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King.

Note added

After this paper was completed, Ref. [34] appeared in which the diphoton excess was discussed in a model with two Higgs doublet fields and a real inert septet field.

Appendix A: One-loop beta functions

In this Appendix, we give the analytic formulae for the beta functions at the one-loop level which are used for the RGE analysis given in Sec. III. The beta functions are defined as

$$\beta(\lambda_i) \equiv \frac{d}{d \ln \mu} \lambda_i, \quad (\text{A1})$$

where μ is an energy scale. For the Yukawa couplings, we only keep the contribution of the top Yukawa coupling y_t .

For the gauge couplings, the beta functions for the $SU(3)_c$ (g_s), $SU(2)_L$ (g) and $U(1)_Y$ (g_Y) coupling are given by

$$\beta(g_s) = -\frac{1}{16\pi^2} 7g_s^3, \quad (\text{A2})$$

$$\beta(g) = \frac{1}{16\pi^2} \left[-\frac{19}{6} + \frac{2}{3}(1^2 + 2^2 + \dots + T^2) \right] g^3, \quad (\text{A3})$$

$$\beta(g_Y) = \frac{1}{16\pi^2} 7g_Y^3, \quad (\text{A4})$$

where $\beta(g_s)$ is the same form as the SM one.

The beta function for y_t is given by

$$\beta(y_t) = \frac{1}{16\pi^2} \left[9y_t^3 - \left(8g_s^2 + \frac{9}{4}g^2 + \frac{17}{12}g_Y^2 \right) y_t \right]. \quad (\text{A5})$$

Finally, we give the beta functions for all the scalar quartic couplings as follows

$$\beta(\lambda_1) = \frac{1}{16\pi^2} \left[12\lambda_1^2 + 4\lambda_3^2 + 2\lambda_4^2 + 4\lambda_3\lambda_4 + 2\lambda_5^2 + 4\rho_1^2(1+2T) + \frac{9}{4}g^4 + \frac{3}{4}g_Y^4 + \frac{3}{2}g_Y^2g^2 - 3\lambda_1(3g^2 + g_Y^2) \right], \quad (\text{A6})$$

$$\beta(\lambda_2) = \frac{1}{16\pi^2} \left[12\lambda_2^2 + 4\lambda_3^2 + 2\lambda_4^2 + 4\lambda_3\lambda_4 + 2\lambda_5^2 + 4\rho_2^2(1+2T) - 12y_t^4 + \frac{9}{4}g^4 + \frac{3}{4}g_Y^4 + \frac{3}{2}g_Y^2g^2 - 3\lambda_2(3g^2 + g_Y^2 - 4y_t^2) \right], \quad (\text{A7})$$

$$\beta(\lambda_3) = \frac{1}{16\pi^2} \left[2(\lambda_1 + \lambda_2)(3\lambda_3 + \lambda_4) + 4\lambda_3^2 + 2\lambda_4^2 + 2\lambda_5^2 + 4\rho_1\rho_2(1+2T) + \frac{9}{4}g^4 + \frac{3}{4}g_Y^4 - \frac{3}{2}g_Y^2g^2 - 3\lambda_3(3g^2 + g_Y^2 - 2y_t^2) \right], \quad (\text{A8})$$

$$\beta(\lambda_4) = \frac{1}{16\pi^2} \left[2\lambda_4(\lambda_1 + \lambda_2 + 4\lambda_3 + 2\lambda_4) + 8\lambda_5^2 + 3g^2g_Y^2 - 3\lambda_4(3g^2 + g_Y^2 - 2y_t^2) \right], \quad (\text{A9})$$

$$\beta(\lambda_5) = \frac{1}{16\pi^2} \left[2\lambda_5(\lambda_1 + \lambda_2 + 4\lambda_3 + 6\lambda_4) - 3\lambda_5(3g^2 + g_Y^2 - 2y_t^2) \right], \quad (\text{A10})$$

$$\beta(\rho_1) = \frac{1}{16\pi^2} \left[\frac{3}{2}T(T+1)g^4 + 6\lambda_1\rho_1 + 4\lambda_3\rho_2 + 2\lambda_4\rho_2 + 8\rho_1^2 + 8\rho_1\lambda_\chi(3+2T) + 4\rho_1\lambda'_\chi \left(6c_0 + \sum_{i=1,T} c_i \right) - \rho_1(6g^2T(T+1) + \frac{9}{2}g^2 + \frac{3}{2}g_Y^2 - 6y_t^2) \right], \quad (\text{A11})$$

$$\beta(\rho_2) = \frac{1}{16\pi^2} \left[\frac{3}{2}T(T+1)g^4 + 6\lambda_2\rho_2 + 4\lambda_3\rho_1 + 2\lambda_4\rho_1 + 8\rho_2^2 + 8\rho_2\lambda_\chi(3+2T) + 4\rho_2\lambda'_\chi \left(6c_0 + \sum_{i=1,T} c_i \right) - \rho_2 \left(6g^2T(T+1) + \frac{9}{2}g^2 + \frac{3}{2}g_Y^2 - 6y_t^2 \right) \right], \quad (\text{A12})$$

$$\beta(\lambda_\chi) = \frac{1}{16\pi^2} \left[\frac{3}{4}T^2(T+1)^2g^4 + 2(\rho_1^2 + \rho_2^2) + 72(\lambda_\chi + c_0\lambda'_\chi)^2 + \sum_{i=1,T} (4\lambda_\chi + c_i\lambda'_\chi)^2 - 12g^2T(T+1)(\lambda_\chi + c_0\lambda'_\chi) \right] - c_0\beta(\lambda'_\chi), \quad (\text{A13})$$

where

$$(c_0, c_1, c_2) = (0, 0, 0) \text{ for } T = 2, \quad (\text{A14})$$

$$(c_0, c_1, c_2, c_3) = (72, 288, 432, 72) \text{ for } T = 3, \quad (\text{A15})$$

$$(c_0, c_1, c_2, c_3, c_4) = (200, 800, 1360, 440, 160) \text{ for } T = 4. \quad (\text{A16})$$

The beta function for λ'_χ is respectively given for the $T = 3$ and $T = 4$ case as

$$\beta(\lambda'_\chi) = \frac{1}{16\pi^2} \left(6120\lambda_\chi'^2 + 96\lambda_\chi\lambda'_\chi + \frac{3}{2}g^4 - 144g^2\lambda'_\chi \right), \quad \text{for } T = 3, \quad (\text{A17})$$

$$\beta(\lambda'_\chi) = \frac{1}{16\pi^2} \left(21992\lambda_\chi'^2 + 96\lambda_\chi\lambda'_\chi + \frac{3}{2}g^4 - 240g^2\lambda'_\chi \right), \quad \text{for } T = 4. \quad (\text{A18})$$

We have checked the consistency of the above formulae with those given in Ref. [17].

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